

PERTURBATIVE HOT GAUGE THEORIES - RECENT RESULTS

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ABSTRACT

Current results in high temperature gauge theories obtained in the context of the perturbative method of resumming hard thermal loops are reviewed. Beyond leading order properties of the gluon excitation, and the recent (controversial) calculations of the damping rates are discussed. QCD predictions on plasma signatures are exemplified by the thermal production rates of energetic as well as soft photons.

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Current results in high temperature gauge theories obtained in the context of the perturbative method of resumming hard thermal loops are reviewed. Beyond leading order properties of the gluon excitation, and the recent (controversial) calculations of the damping rates are discussed. QCD predictions on plasma signatures are exemplified by the thermal production rates of energetic as well as soft photons.

1. Introduction

At this Workshop many talks have been devoted to the properties of systems at high temperatures and densities. In the following I will concentrate on hot gauge theories and try to summarize some of the recent developments and open questions in the framework of the hard thermal loop (HTL) resummation method^{1,2,3,4}. I mainly follow the advice by Braaten and Pisarski that "in the end, the best way to demonstrate the consistency of the effective expansion is to *apply* it to a wide range of *physical processes*".

New theoretical results to be shortly discussed are on:

- the gluon excitation beyond leading order;
- damping rates for the excitations at rest as well as the energetic ones;
- the dynamical screening mechanism of mass singularities;
- the production of soft photons (with energies of $O(gT)$) from a quark gluon plasma and problems related to the evaluation of their rate.

2. Gluon self-energy: next-to-leading order corrections

2.1. Gluon plasma frequency

By an impressive calculation H. Schulz⁵ succeeded to determine beyond leading order the real part of the (longitudinal) gluon self-energy in the long-wave length

limit. The corresponding solution of the dispersion relation for the gluon excitation at rest,

$$\Omega^2 - \Pi_L(\Omega, \vec{q} = 0) = 0, \quad (1)$$

becomes (with $N_f = 0$ quarks and for N colours) with $w \equiv Re\Omega$ ($\equiv m_g$),

$$w^2 = \frac{g^2 N}{9} T^2 (1 - 0.18 g \sqrt{N} + \dots), \quad (2)$$

where the corrections to the well known leading order HTL term⁶ ($m_g \simeq g\sqrt{N}T/3$) arise from (i) hard two-loop and (ii) hard and effective soft one-loop diagrams - after the zeroth order HTL's are subtracted.

Contrary to the original estimate² based on power counting arguments Schulz's explicit evaluation of 13 two-loop diagrams (i) with hard internal momenta shows that they do not contribute to Eq.(2) at $O(g^3 T^2)$, but only at higher order $O(g^4 T^2)$! A clever rearrangement of the one-loop terms (ii) simplifies their treatment. One sample of terms - the "one-loop hard" ones in the notation of⁵ - gives a negligible contribution of $O(g^4 \ln 1/g T^2)$: this, however, can only be seen after detailed calculations as it is proportional to $g^4 N T^2 I(g\sqrt{N})$, in terms of a principal value integral

$$I(g\sqrt{N}) = \frac{1}{4\pi^2} \mathbf{P} \int_0^\infty \frac{x dx}{x^2 - 1} \frac{1}{\exp(\frac{g\sqrt{N}}{6} x) - 1}. \quad (3)$$

After expanding the Bose-Einstein factor for $g \rightarrow 0$ one expects $I(g\sqrt{N}) \simeq 1/g$, i.e. an overall contribution to Eq.(2) of $O(g^3 T^2)$. However, since $\mathbf{P} \int_0^\infty \frac{dx}{x^2 - 1} = 0$, the integral behaves as

$$I(g\sqrt{N}) \underset{g \rightarrow 0}{\simeq} \frac{1}{8\pi^2} \ln g, \quad (4)$$

and this contribution becomes negligible.

Indeed, the $O(g^3 T^2)$ corrections to the plasma frequency, Eq.(2), are only contained in the difference of the dressed minus the bare self-energy contributions (including the tadpole diagrams), with the one-loop momentum restricted to the soft integration region. After a lengthy analytical evaluation the value of the coefficient in Eq.(2) is determined numerically. It has a negative sign, which is familiar from the behaviour of the QCD-pressure ($N_f = 0$)⁷:

$$P \simeq \frac{(N^2 - 1)\pi^2}{45} T^4 (1 - \frac{5g^2 N}{16\pi^2} + \dots), \quad (5)$$

i.e. indicating as Eq.(2) a phase transition for a critical value of $g_{cr} \simeq 3.2$ for $N = 3$.

Schulz' calculation shows that (i) the correction is (covariant) gauge parameter independent in the algebraic sense⁸, i.e. the dependence is weighted by the on-shell condition; (ii) it is also independent of the momentum cut-off between the soft- and hard loop momentum; and (iii) there is consistency with the evaluation of the damping rate in the long-wave length limit via the dressed soft one-loop diagrams⁹.

2.2. QCD - Debye mass

A new analysis at next-to-leading order of the static longitudinal gluon self-energy $\Pi_{00}(w = 0, \vec{q})$ and its relation to the QCD-Debye mass is performed by A. Rebhan¹⁰. Here only those of his results are presented which are of relevance in connection with problems to be discussed in the following sections.

Instead of the usual QED-type definition of the QCD-Debye mass m by $m^2 \triangleq \Pi_{00}(w = 0, \vec{q} \rightarrow 0)$, Rebhan proposes the proper definition of m by the pole position of the longitudinal gluon propagator, i.e. by

$$m^2 \equiv \Pi_{00}(w = 0, \vec{q}^2 = -m^2). \quad (6)$$

In covariant gauges the Braaten-Pisarski resummation gives in the effective one-loop approximation for $\vec{q}^2 \simeq -m^2$

$$\Pi_{00}(0, \vec{q}) \simeq m^2 + \delta\Pi_{00}, \quad (7)$$

where

$$\delta\Pi_{00}/g^3T^2 \simeq c_1 \ln(\vec{q}^2 + m^2) + c_2 + \xi (\vec{q}^2 + m^2) I(\vec{q}, m). \quad (8)$$

In order to obtain a sensible value for the Debye mass m the logarithmic singularity present at the next-to-leading order has to be "screened". Since it is due to the massless transverse gluon propagator in the loop a (nonperturbative) "magnetic mass" m_{mag} may be used as infrared cut-off.

Concerning the gauge parameter, i.e. ξ , dependence one observes in Eq.(8) an algebraic on-shell independence⁸, in case the integral $I(\vec{q}, m)$ is well behaved at $\vec{q}^2 = -m^2$: actually, however, it has a linear on-shell divergence. Therefore gauge independence of m^2 at $O(g^3)$ requires infrared regularisation before the on-shell limit is performed¹¹: e.g. regularisation in $D = 3 + 2\epsilon, \epsilon > 0$, dimensions leads to a ξ -dependence

$$\delta\Pi_{00} \simeq g^3T^2\xi (\vec{q}^2 + m^2)^{2\epsilon}, \quad (9)$$

and gauge parameter independence then follows on-shell when keeping $\epsilon > 0$.

In summary the QCD-Debye mass becomes

$$m^2 \simeq \frac{g^2N}{3}T^2\left(1 + \frac{\sqrt{3N}}{2\pi}g \ln\left(\frac{\text{const}}{g}\right) + \dots\right), \quad (10)$$

when $m_{mag} \simeq g^2T$ and $N_f = 0$ is assumed.

3. Damping rates

3.1. Gluon and quark excitations at rest

For the understanding of the QCD plasma properties at high temperature T it is necessary to calculate the damping rates of the plasma excitations (waves), the

bosonic as well as the fermionic ones. First I concentrate on the damping rates at zero momentum. In Coulomb gauge the important results are:
for the (spatially transverse and longitudinal) gluon excitation⁹

$$\gamma_{T,L} \simeq 6.63 \frac{g^2 NT}{24\pi} , \quad (11)$$

i.e. a positive constant implying the resolution of the long standing "plasmon problem" ;

for the quark excitation^{12,13} (with positive and negative helicity)

$$\gamma_{\pm} \simeq 5.71 \frac{g^2 C_f T}{16\pi} , \quad \text{for } N_f = 3 . \quad (12)$$

C_f is the Casimir for the fundamental $SU(N)$ representation.

When calculating these damping constants in arbitrary covariant gauges, one has to be very careful^{14,15,16}. In order to determine the gauge fixing, i.e. ξ , dependence of the resummed gluon and quark self energies, which only comes from the gauge dependent terms in the resummed gluon propagator, it is convenient to use the Ward identities¹⁷ for the effective propagators and vertices. It is found that the gauge variations are proportional to the corresponding inverse propagators. E.g. the gauge variation of the imaginary part of the transverse gluon self energy is:

$$Im \Delta^* \Pi_T(\omega) = -\xi g^2 N/2 q_T^4 I_T(q) + \dots \quad (13)$$

where the mass-shell is defined by

$$q_T^2 \underset{\omega \rightarrow m_g}{\simeq} m_g (m_g - \omega) . \quad (14)$$

On shell $Im \Delta^* \Pi_T$ vanishes provided that the corresponding integral

$$I_T(q) = Im \sum_{p_0} \int_{\text{soft}} \frac{d^3 p}{k^4 p_T^2} , \quad q = p + k, \quad (15)$$

does not develop poles on the mass-shell⁸. However, it is singular near the mass-shell,

$$I_T(\omega) \simeq \frac{1}{16\pi} \frac{T}{m_g} \frac{1}{(\omega - m_g)^2} , \quad (16)$$

which implies gauge dependence for the damping rate at rest¹⁵ (cf. Eq. (11)):

$$\delta\gamma_T \propto \xi g^2 NT . \quad (17)$$

An analogous result holds for the quark rate¹⁶.

Eq. (16) indicates that this gauge dependence is related to infrared, i.e. to mass singularities, despite the fact that the damping rate itself is finite.

As already discussed in subsection 2.2 this problem is resolved by interchanging limits¹¹, i.e. following the prescription of keeping the infrared regulator - e.g. ϵ in dimensional regularisation - different from zero when taking the on mass-shell limit.

Therefore for calculating damping rates at finite T an infrared regulator is necessary, at least in covariant gauges¹⁸. This may be a hint for a rather complicated - gauge dependent - singularity structure of propagators at finite T , especially near the mass-shell¹⁹.

3.2. Fast moving excitations

Recent studies of damping rates γ of fast moving particles in hot QED or QCD do not yet offer satisfactory results.

As the simplest case I discuss the damping rate of a heavy fermion of mass M . The energy, the momentum and the velocity of the fermion are denoted by E, \vec{p} and v , respectively. The main interest is the leading order behavior, i.e. for $g \rightarrow 0$ at high temperature T with $M > T$ in the limit $v \rightarrow 1$.

The difficulties arise from the infrared sensitive behaviour of the rate, which in a first approximation is related to the Rutherford cross section and the fermion density $n(T)$:

$$\gamma \approx n(T) \sigma_{Rutherford} \approx T^3 g^4 \int \frac{dq}{q^3}, \quad (18)$$

i.e. showing a quadratic divergence. With an infrared cut-off of $O(gT)$ γ becomes “anomalous”: its magnitude is proportional to $g^2 T$! Including the HTL resummation the damping rate - in the one-loop approximation, in which the hard energetic fermion/quark emits/absorbs one soft (dressed) boson/gluon with momentum q - may be expressed as:

$$\gamma(p_0) \simeq g^2 C_f T \int_{\text{soft}} \frac{d^3 q}{(2\pi)^3} \int_{-q}^{+q} \frac{dq_0}{q_0} \rho_T(q_0, q) \text{Im} G^R(p_0 - q_0, \vec{p} - \vec{q}), \quad (19)$$

replacing and generalising Eq. (18). The dominant transverse spectral density $\rho_T(q_0, q)$ enters, with the behaviour for $q \rightarrow 0$:

$$\int_{-q}^{+q} \frac{dq_0}{q_0} \rho_t(q_0, q) \simeq 1/q^2. \quad (20)$$

The heavy fermion propagator is approximated by the (retarded) function G^R : inserting G^R of a free fermion one finds - instead of Eq. (18) - a logarithmically divergent rate:

$$\gamma(p^0 \simeq E) \simeq \frac{g^2}{4\pi} C_f T \int^{gT} \frac{dq}{q}. \quad (21)$$

In QCD this integral becomes finite by introducing $m_{\text{mag}} \simeq g^2 T$: evaluating γ on the real axis ($p^0 \simeq E$) a finite value may be - and has been - obtained for quarks^{20,21,22,23}, which is even gauge parameter independent²⁴.

In hot QED this logarithmic divergence cannot be cutoff by the magnetic mass. Therefore a self-consistent determination of the damping rate may be attempted as it has been first conjectured by Lebedev and Smilga²⁵: γ itself is the infrared cut-off^{26,27}. The original proposal amounts to take for G^R in Eq. (19) a Lorentzian ansatz,

$$G^R \simeq \frac{1}{p_0 - (E(p) - i\gamma)} . \quad (22)$$

However, clarification of the “on-shell” condition is required, i.e. either one keeps p^0 on the real axis, $p^0 \simeq E(p)$, or one demands self-consistency to hold on the complex pole $p^0 = E(p) - i\gamma$. Under the assumption that G^R is given by Eq. (22) with a complex pole on the first (and only) sheet in the energy plane, it turns out²⁷ that the infrared divergence is not screened by a non-vanishing γ ! Therefore, this attempt especially fails for QED when Eq. (22) is used for the “dissipative” retarded fermion propagator.

For the fast QCD damping rates Pisarski²⁸ tried the self-consistent approach on the complex pole, but introducing in addition the magnetic mass as an infrared cut-off, i.e. replacing the r.h.s. of Eq. (20) by $1/(q^2 + m_{mag}^2)$. Under the condition $m_{mag} \gg \gamma$, which separates the pole from a nearby branch point, his result is (for $v = 1$) :

$$\gamma \simeq \frac{g^2}{8\pi} C_f T \ln\left(\frac{m_g^2}{m_g^2 + 2\gamma m_{mag}}\right) . \quad (23)$$

However, it has been pointed out²⁹ that this result is not stable: assuming instead of a constant γ an energy dependent one, no self-consistent solution of Eqs. (19) and (22) is found for the case $v < 1$, whereas for $v \simeq 1$ the solution is - instead of Eq. (23) :

$$\gamma(E) \simeq \frac{g^2}{8\pi} C_f T \left(\frac{6\pi m_g^2}{Em_{mag}} \right) . \quad (24)$$

Because of the described difficulties and inconsistencies of these results, it is even concluded that the “anomalous” damping of fast quanta is not observable, and thus unphysical³⁰. However, because of the close (?) connection of γ to transport phenomena^{19,23}, especially to colour diffusion³¹, one has to argue that the simple ansatz, Eq. (22), for the retarded energetic fermion propagator G^R is not reflecting realistic physical conditions. From general properties of the spectral functions, or equivalently from the cutting rules at finite temperature³², one may deduce that retarded Green functions do not have complex poles on the “physical” sheet^{28,33}. An explicit and simple example³⁴, in which the pole is present only on unphysical sheets, shows that self-consistency for γ may be possible without introducing m_{mag} . In QED a fermion damping rate at leading order

$$\gamma(p^0 \simeq E(p)) \simeq \frac{e^2}{4\pi} T \ln \frac{eT}{\gamma} \simeq \frac{e^2}{4\pi} T \ln \frac{1}{e} \quad (25)$$

is then derived.

4. Production of photons from a quark - gluon plasma

Information about the properties of the QCD plasma in its initial stage of formation in heavy – ion collisions is expected to be provided by photons, real as well as virtual ones^{35,36}.

At high temperatures the emission of hard real photons is determined at lowest order in the electromagnetic coupling and in g by the basic QCD processes: quark-antiquark annihilation ($q\bar{q} \rightarrow \gamma g$) and Compton scattering ($qg \rightarrow \gamma q$). However, at this order the thermal production rates are logarithmically divergent due to massless quark exchange. This is in contrast to the rate of thermal heavy photons, where the quark mass singularities cancel. Therefore, for real photons the mass singularities have to be shielded by thermal effects in order to derive infrared safe predictions.

Indeed HTL resummation provides finite rates for *hard* photons because of Landau damping on the exchanged quark: one considers an effective one-loop diagram with a quark loop with one soft (dressed at $O(gT)$) and one hard (bare) quark line. No effective vertices have to be included.

The resulting emission rate for real photons with energy E is^{37,38} :

$$E \frac{dW^\gamma}{d^3p} \simeq \frac{e_q^2 \alpha \alpha_s}{2\pi^2} T^2 e^{-E/T} \ln \left(\frac{c}{\alpha_s} \frac{E}{T} \right), \quad (26)$$

where even the constant c under the logarithm is determined, $c \simeq 0.23$, after the hard annihilation and Compton contributions are added. The emission rate is seen to be independent of the cut-off between soft and hard momenta. It contains, however, a logarithmic dependence on the strong coupling constant, but the mass singularities are dynamically screened.

Other successful examples are e.g. the production of photon pairs³⁹ and hard lepton pairs⁴⁰, or processes (e.g. responsible for the collisional energy loss) in which massless gauge bosons^{41,42,43}, photons or gluons, are exchanged: dynamical screening indeed cures logarithmic mass divergencies.

For *soft* external photons with energies E of $O(gT)$, or softer, additional screening processes - similar to the dilepton case^{44,45} - have to be included into the theoretical description^{46,47}: evaluating the dominant contribution to the production rate both internal quark propagators are assumed to be soft, and the vertices have also to be resummed. As the internal quark propagators are resummed no divergence appears when the quark momenta are vanishing. However, the introduction of effective HTL vertices leads to unscreened collinear divergences. This may be seen from the effective quark - photon vertex^{2,3} :

$$^*\Gamma^\mu = \gamma^\mu + m_f^2 \int \frac{d\Omega}{4\pi} \frac{Q^\mu \not{Q}}{(Q \cdot k)(Q \cdot k')}, \quad (27)$$

where the second term is the HTL correction – in terms of an angular integral – to the bare vertex γ^μ . Q is a light-like vector, $Q^2 = 0$, and $m_f^2 = 2\pi\alpha_s T^2/3$. The mass

singularity arises from the $1/(Q \cdot k)$ factors when both quark momenta, k and k' , are space-like, i.e. from:

$$\mathbf{P} \left(\frac{1}{Q \cdot k} \right) \delta(Q \cdot k') = \left(\frac{1}{Q \cdot p} \right) \delta(Q \cdot k') , \quad (28)$$

when $(Q \cdot p) = 0$ for the light-like photon momentum $p^2 = 0$. Regularising this singularity by using dimensional regularisation of the angular integral over $d\Omega$ in $D = 3 + 2\epsilon$ dimensions the leading (singular) contribution to the soft photon production rate for $g \rightarrow 0$ then reads⁴⁶ :

$$E \frac{dW^\gamma}{d^3\vec{p}} \simeq \frac{1}{\epsilon} \frac{e_q^2 \alpha \alpha_s}{2\pi^2} T^2 n_F(E) \left(\frac{m_f}{E} \right)^2 \ln \left(\frac{1}{\alpha_s} \right) . \quad (29)$$

This result shows that the Braaten-Pisarski resummation does not yield a finite soft real photon production rate: a logarithmic divergence remains. At present we do not know how to screen this mass singularity by a consistent procedure. Eq. (29) is valid for soft massless, i.e. non-thermalized photons: the quark - gluon plasma has to have a finite size, such that its characteristic length is smaller than the photon's mean free path.

One may identify the diagrams which are responsible for the singularities. The massless quark exchange shows up in two \rightarrow three amplitudes, e.g. for $g^* q \rightarrow q^* g\gamma$ (* denotes dressed partons), and the singularity arises from the configuration $Q \cdot p = 0$ (cf. Eq. (28)), corresponding to a collinear singularity; here Q is identified as the momentum carried by the bare gluon in the HTL vertex.

Important processes, which are sensitive to scales $\leq O(g^2T)$, and therefore have to be described beyond the HTL resummation, e.g. photon bremsstrahlung from a QED plasma⁴⁸, or from a QGP^{49,50} by including the effect of Landau-Pomeranchuk suppression⁵¹, and the energy loss due to radiation^{52,53,54} of quarks and gluons traversing in the plasma are under detailed studies. However, more work is needed in order to better control the presently applied approximations.

5. Conclusion

During the last years an impressive amount of work related to the resummation method of hard thermal loops has been performed⁵⁵, which improves our understanding of the behaviour of gauge theories at high temperatures. For scales from $O(gT)$ to T the resummation turns out to be in general successful in screening fermionic and gluonic mass singularities, i.e. infrared stable predictions are possible at leading order. One exception, however, is the rate for the production of real photons with energies of $O(gT)$.

Still unsolved problems remain with respect to phenomena on scales of $O(g^2T)$ and smaller. The evaluation of the damping rates of the (energetic) plasma excitations illustrates the present difficulties, including the question of the existence of the magnetic screening mass. Concerning the physical significance of these rates

a detailed description is necessary, e.g. how hard quark and gluon jets produced during the very early stages (on time scales $\simeq 0.1fm$) in heavy-ion collisions may probe the dense QCD plasma subsequently formed at times of $O(1fm)$.

The fact that the QCD coupling constant is rather large, $g \simeq 1$, at temperatures of the order of a few hundreds of MeV, which is the expected realistic range for studying experimentally quark-gluon plasma properties, requires large extrapolations of the discussed perturbative results. So far they are derived predominantly in leading, at most next-to-leading order for high T and small g .

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